Asymmetric nuclear matter in a Hartree-Fock approach to non-linear QHD

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Abstract

The Equation of State (EOS) for asymmetric nuclear matter is discussed starting from a phenomenological hadronic field theory of Serot-Walecka type including exchange terms. In a model with self interactions of the scalar sigma-meson we show that the Fock terms naturally lead to isospin effects in the nuclear EOS. These effects are quite large and dominate over the contribution due to isovector mesons. We obtain a potential symmetry term of "stiff" type, i.e. increasing with baryon density and an interesting behaviour of neutron/proton effective masses of relevance for transport properties of asymmetric dense matter.

Phenomenological hadronic field theories (Quantum Hadrodynamics, QHD) are widely used in dense nuclear matter studies since relativistic effects are expected to increase with baryon density [1]. In most of the previous works on the subject, the Relativistic Mean Field (RMF) approximation of QHD has been followed. In the RMF the nucleon and meson fields are treated as classical fields and consequently a Hartree reduction of one body density matrices is used.

Although the model has driven a large amount of results on relativistic effects in nuclear structure and dynamics [2–6], the lack of exchange terms has implied some non satisfying features of the theory and some efforts have been done to try to cure this problem [7–10]. In the RMF theory each meson field is introduced, with appropriated readjusted couplings, just to describe the dynamics of a corresponding degree of freedom, without mixing due to many-body effects. Neutral σ and ω mesons are in charge of saturation properties, isospin effects are carried by isovector δ [$a_0(980)$] and ρ mesons and finally spin terms are coming from pseudoscalar π and η fields. In a sense the model represents a straightforward extension of the One-Boson-Exchange (OBE) description of nucleon-nucleon scattering.

The aim of this letter is to introduce explicit many-body effects just evaluating exchange term contributions. As we already know from non-relativistic effective interactions, like the Skyrme forces, Fock terms play an essential role in symmetry breakings and consequent mixing of different degrees of freedom. Similar effects are expected here. In particular, in the context of the QHD model, essential properties of nuclear matter come mostly from the two neutral strong meson fields. Hence it is important to evaluate the Fock contribution associated with these fields.

We will focus our attention on isospin contributions to the nuclear EOS, symmetry term and neutron/proton effective masses, and on the relativistic transport equation for asymmetric nuclear matter.

Then we will start from a $QHD-II$ model [1] where the nucleons are coupled to neutral scalar σ and vector ω mesons and to the isovector ρ meson. Self-interaction terms of the σ-field were originally introduced for renormalization reasons [11,12] and can also be considered as a way to parametrize the density dependence of the NN force. Actually they are also describing medium effects essential to reproduce important properties (compressibility and nucleon effective mass) of nuclear matter around saturation density. The Lagrangian density for this model is given by:

$$\mathcal{L} = \bar{\psi}[\gamma_\mu(i\partial^\mu - g_V\mathcal{V}^\mu - g_\rho\mathcal{B}^\mu \cdot \tau) - (M - g_S\phi)]\psi + \frac{1}{2}(\partial_\mu\phi\partial^\mu\phi - {m_S}^2\phi^2)$$

$$-\frac{a}{3}\phi^3 - \frac{b}{4}\phi^4 - \frac{1}{4}W_{\mu\nu}W^{\mu\nu} + \frac{1}{2}{m_V}^2\mathcal{V}_\nu\mathcal{V}^\nu - \frac{1}{4}\mathbf{L}_{\mu\nu}\cdot\mathbf{L}^{\mu\nu} + \frac{1}{2}{m_\rho}^2\mathcal{B}_\nu\cdot\mathcal{B}^\nu \tag{1}$$

where $W^{\mu\nu}(x) = \partial^\mu\mathcal{V}^\nu(x) - \partial^\nu\mathcal{V}^\mu(x)$ and $\mathbf{L}^{\mu\nu}(x) = \partial^\mu\mathcal{B}^\nu(x) - \partial^\nu\mathcal{B}^\mu(x)$.

Here $\psi(x)$ generally denotes the fermionic field, $\phi(x)$ and $\mathcal{V}^\nu(x)$ represent neutral scalar and vector boson fields, respectively. $\mathcal{B}^\nu(x)$ is the charged vector field and τ denotes the isospin matrices.

From the previous equation one can derive the field equations and the canonical energy-momentum tensor [1].

In our approach we will perform the many-body calculations in the quantum phase space introducing the Wigner transform of the one-body density matrix of the fermion field. This method has two main advantages, the use of physical quantities and the direct derivation of dynamical transport equations. The one–particle Wigner function is defined as:

$$[\widehat{F}(x,p)]_{\alpha\beta} = \frac{1}{(2\pi)^4}\int d^4R e^{-ip\cdot R} <: \bar{\psi}_\beta(x+\frac{R}{2})\psi_\alpha(x-\frac{R}{2}) :> \quad ,$$

here α and β are indices for intrinsic degrees of freedom of the fermionic field (spin and isospin). The brackets denote statistical averaging and the double dots denote normal ordering. The equation of motion for the Wigner function can be derived from the Dirac field equation by using standard procedures (see e.g. Refs. [13,14]), it reads:

$$\begin{aligned}
&\frac{i}{2}\partial_\mu[\gamma^\mu\hat{F}(x,p)]_{\alpha\beta} + p_\mu[\gamma^\mu\hat{F}(x,p)]_{\alpha\beta} - M\hat{F}_{\alpha\beta}(x,p) \\
&-g_V\frac{1}{(2\pi)^4}\int d^4R e^{-ip\cdot R} <: \bar{\psi}_\beta(x_+)\gamma^\mu_{\alpha\delta}\psi_\delta(x_-)\mathcal{V}_\mu(x_-) :> \\
&+g_S\frac{1}{(2\pi)^4}\int d^4R e^{-ip\cdot R} <: \bar{\psi}_\beta(x_+)\psi_\alpha(x_-)\phi(x_-) :> \\
&-g_\rho\frac{1}{(2\pi)^4}\int d^4R e^{-ip\cdot R} <: \bar{\psi}_\beta(x_+)\gamma^\mu_{\alpha\delta}\psi_\delta(x_-)\tau\cdot\mathcal{B}_\mu(x_-) :>= 0 \ ,
\end{aligned} \tag{2}$$

with $x_+ = x + \frac{R}{2}$ and $x_- = x - \frac{R}{2}$.

In order to take into account the contribution of exchange terms in a manageable way we assume, as a basic approximation, that in the equations of motion for the meson fields the terms containing derivatives can be neglected with respect to the mass terms. Therefore the meson field operators are directly connected to the operators of the nucleon scalar and current densities:

$$\begin{aligned}
\widehat{\Phi}/f_S + A\widehat{\Phi^2} + B\widehat{\Phi^3} &= \bar{\psi}(x)\psi(x) \ , \\
\widehat{V}^\mu(x) &= f_V\bar{\psi}(x)\gamma^\mu\psi(x) \ , \\
\widehat{\mathbf{B}}^\mu(x) &= f_\rho\bar{\psi}(x)\gamma^\mu\tau\psi(x) \ , \qquad (3)
\end{aligned}$$

where $f_S = (g_S/m_S)^2$, $f_V = (g_V/m_V)^2$, $f_\rho = (g_\rho/m_\rho)^2$ and $\widehat{\Phi} = g_S\phi$, $\widehat{V}^\mu = g_V\mathcal{V}^\mu$, $A = a/g_S^3$, $B = b/g_S^4$, $\mathbf{B}^\mu = g_\rho\mathcal{B}^\mu$. After substituting in Eq.(2) these expressions for the meson field operators, we obtain an equation which contains only nucleon field operators.

The present approximation implies that retardation and finite size effects in the exchange of mesons between nucleons are neglected. However, we are concerned with a semiclassical description of nuclear dynamics, so that the nuclear medium is supposed to be in states for which the nucleon scalar and current densities are smooth functions of the space-time coordinates. Therefore, because of the small Compton wave–lengths of the mesons σ, ω and ρ, the assumptions expressed by Eq.s(3) are quite reasonable.

An attempt to include exchange terms in the QHD approach was previously performed without self–interaction terms for the σ field [9], with results not satisfying due to the inadequacy of the model. Within the same model, a relativistic kinetic equation with self–consistent mean field has been derived in ref. [15] taking into account the exchange terms. Here we evaluate the effects in a more physical model with self-interacting higher order σ terms.

In order to introduce the Fock terms, we consider the expansion of the field operator $\widehat{\Phi}$ in terms of $\widehat{\rho_S} = \bar{\psi}\psi$: $\widehat{\Phi} = \sum_n \chi_n \widehat{\rho_S}^n$. Then, the statistical average of $\widehat{\rho_S}^n$ can be written, at the Hartree-Fock level as it follows:

$$<: \widehat{\rho_S}^n(x) :>= \rho_S^n(x) - \frac{1}{k!}\sum_{k=2}(-1)^k Tr\widehat{F}^k(x)\frac{d^{k+1}\rho_S^n(x)}{d\rho_S^{k+1}(x)}, \qquad (4)$$

where $\rho_S(x) =<: \bar{\psi}(x)\psi(x) :>$ and $\widehat{F}(x) = \int d^4p \ \widehat{F}(x,p)$.

The statistical averages of the products that appear in Eq.(2) can be expressed, by considering the expansion in Eq.(4) up to the next-to-leading term, as shown in ref. [16]. In particular the matrix $\widehat{F}(x)$ can be decomposed in components with definite transformation properties (Clifford algebra), where, for instance, the scalar and the vector parts (F,F_μ) are related to scalar and current densities as: $8\ F(x) = \rho_S(x)$, $8\ F^\mu(x) = j^\mu(x) =<: \bar{\psi}(x)\gamma^\mu\psi(x) :>$. Therefore the coefficients of the expansion Eq.(4) contain a factor that decreases as $(1/8)^k$.

In order to evaluate the nuclear Equation of State the quantity of interest is the statistical average of the canonical energy-momentum density tensor. According to our approximation, where terms containing the derivatives of the meson fields are neglected, the tensor takes the form:

$$T_{\mu\nu}(x) = \frac{i}{2}\bar{\psi}(x)\gamma_\mu\partial_\nu\psi(x) + [U(\widehat{\Phi}) - \frac{1}{2}\widehat{V}_\lambda\widehat{V}^\lambda/f_V - \frac{1}{2}\widehat{\mathbf{B}}_\lambda\widehat{\mathbf{B}}^\lambda/f_\rho]g_{\mu\nu}, \tag{5}$$

where $g_{\mu\nu}$ is the diagonal metric tensor and $U(\widehat{\Phi}) = \frac{1}{2}\widehat{\Phi}^2/f_S + A/3\ \widehat{\Phi}^3 + B/4\ \widehat{\Phi}^4$. According to the approximation introduced in Ref. [16], the energy-momentum density tensor is given by:

$$\begin{aligned} < T_{\mu\nu}(x) >= 8\int d^4p\ p_\nu F_\mu + \{U(\Phi) - f_V/2\ J_\lambda J^\lambda\ - f_\rho/2\ b_\lambda b^\lambda\}g_{\mu\nu} \\ -\frac{1}{2}[Tr\widehat{F}^2\frac{d^2U(\Phi)}{d\rho_S^2} - f_V Tr(\gamma_\lambda\widehat{F}\gamma^\lambda\widehat{F}) - f_\rho Tr(\gamma_\lambda\tau\widehat{F}\gamma^\lambda\tau\widehat{F})]g_{\mu\nu} \end{aligned} \tag{6}$$

The quantities in square brackets are the contributions of the exchange terms. It is essential to note that Fock terms contain traces of powers of $\widehat{F}(x)$ that naturally bring scalar, vector, tensor, pseudoscalar and pseudovector contributions. In particular for the case of asymmetric nuclear matter we obtain scalar and vector isovector contributions to the EOS, generally associated respectively with δ and ρ mesons. From Eq.(6) we obtain the energy density for asymmetric nuclear matter that in analogy to the Hartree case can be rewritten in the following form:

$$\epsilon = < T_{00} > = {\epsilon_{kin}}^p + {\epsilon_{kin}}^n + U(\Phi) + \frac{1}{2}\tilde{f}_S{\rho_S}^2 + \frac{1}{2}\tilde{f}_V{\rho_B}^2 + \frac{1}{2}f'_S b^2 + \frac{1}{2}f'_V {b_0}^2 \tag{7}$$

where ρ_S, ρ_B are the scalar and baryon densities and $b = \rho_{Sn} - \rho_{Sp}$, $b_0 = \rho_{Bn} - \rho_{Bp}$ are the corresponding isovector parts. The

$$\epsilon_{kin}{}^i = \frac{2}{(2\pi)^3} \int d^3p \sqrt{p^2 + M_i{}^{*2}} = \frac{1}{4}[3E_i{}^* \rho_{Bi} + M_i{}^* \rho_{Bi}] \quad i = n, p \tag{8}$$

are kinetic contributions and

$$\tilde{f}_S = \frac{1}{2} f_V - \frac{1}{8}(\frac{d\Phi}{d\rho_S} + \rho_S \frac{d^2\Phi}{d\rho_S^2}) + \frac{3}{2}\ f_\rho;$$

$$f'_S = \frac{1}{2} f_V - \frac{1}{8}(\frac{d\Phi}{d\rho_S} + \rho_S \frac{d^2\Phi}{d\rho_S^2}) - \frac{1}{2}\ f_\rho;$$

$$\tilde{f}_V = \frac{5}{4} f_V + \frac{1}{8}(\frac{d\Phi}{d\rho_S} - \rho_S \frac{d^2\Phi}{d\rho_S^2}) + \frac{3}{4}\ f_\rho;$$

$$f'{}_V = \frac{1}{4} f_V + \frac{1}{8}(\frac{d\Phi}{d\rho_S} - \rho_S \frac{d^2\Phi}{d\rho_S^2}) + \frac{3}{4}\ f_\rho \tag{9}$$

are density dependent effective coupling constants. $E_i^* = \sqrt{|\mathbf{p}|^2 + M_i^{*2}}$ and M_i^* are the effective masses, see Eq.s.(12,13).

Here we explicitly obtain a density dependence arising also in the vector, isovector and isoscalar couplings, like in the phenomenological approach of ref. [17]. We have verified that the used approximation leads to a thermodynamically consistent theory [16].

We remark that, as in non-linear mean-field models, we have in total five parameters. As usual, the ones related to isoscalar mesons, f_V, f_S, A, B are fixed in order to reproduce equilibrium properties of symmetric nuclear matter: saturation density $\rho_0 = 0.16\ fm^{-3}$, binding energy $E/A = -16\ MeV$, nucleon effective (or Dirac) mass at ρ_0, $M_0{}^* = 0.73\ M$ (see Eq.(13)) and compressibility modulus $K_0 = 245\ MeV$. The coupling constant f_ρ can then be adjusted in order to get a good value for the symmetry energy at saturation density, but now taking into account the contribution to the isovector channel coming from the isoscalar mesons through the Fock terms. In our calculations we have a symmetry coefficient of the Weiszaecker mass formula $a_4 = 31.5\ MeV$. In the Table we list the values obtained for the five parameters.

f_V (fm^2)	f_S (fm^2)	A (fm^{-1})	B	f_ρ (fm^2)
3.998	9.731	0.088	-0.015	0.6

Table $NLHF$ parameters from the fit to saturation properties of nuclear matter (see the text).

According to these values, we see that the dominant contributions to the density dependent coupling functions $\tilde{f}_s$, $\tilde{f}_v$, $\tilde{f}'_s$, $\tilde{f}'_v$, Eq.(9), come essentialy from the isoscalar σ and ω mesons.

We discuss now some results for the EOS of asymmetric nuclear matter. We show the comparison between our *Non Linear Hartree-Fock* ($NLHF$) present calculations and those of the *Non Linear Hartree* (NLH) model of ref. [1,12], including the isovector ρ and δ mesons [18], with parameters chosen in order to give the same saturation properties.

The comparison for symmetry energies per nucleon is presented in Fig.1. As already stressed, in the $NLHF$ results a large contribution to the symmetry term comes from the Fock terms associated with the σ and ω mesons, with the corresponding four parameters of the theory fitted on properties of symmetric NM. Therefore the inclusion of the Fock terms (solid line, $NLHF$) can give the correct value of a_4 even with a relatively small coupling constant for the ρ meson: $f_\rho = 0.6\ fm^2$, close to the free space value $f_{\rho,free}$, see ref. [1].

We show also NLH calculations including both ρ and the isovector scalar δ mesons (dashed line). In this case the coupling constants f_ρ and f_δ have been chosen in order to reproduce at saturation density the same symmetry energy and neutron-proton effective mass splitting that we get within our model. We remark that now we need a $f_\rho = 2.3\ fm^2$, *about four times the free space value*, and also a relatively strong δ coupling, $f_\delta = 1.4\ fm^2$, but still in the range of free space values [19,18]. For reference we show also the result of a NLH calculations including *only* the ρ contribution (long-dashed line). In order to have the same a_4 value at saturation density we need a $f_\rho = 1.2\ fm^2$, still almost two times the free space value. We remark that the inclusion of the δ contribution in the Hartree scheme, necessary for the neutron-proton mass splitting, gives also an attractive term in the symmetry energy, see ref. [18], and so a much stronger ρ coupling is required in order to reproduce the correct a_4 coefficient around saturation density.

In all these relativistic models a quite repulsive density dependence of the symmetry term of the EOS is obtained.

We notice that the density dependence of the symmetry energy that one obtains in the complete $NLH + \rho + \delta$ model is quite different with respect to our results. This is due to the fact that in the $NLHF$ model the coupling functions in the isovector channels (f'_S, f'_V) become density dependent. Such density dependence is shown in Figure 2. In particular we stress the behaviour at sub-nuclear densities due to the opposite sign of the $d\phi/d\rho_S$ contribution, see Eq(9). This implies a "softer" behaviour of the potential symmetry term *below saturation density* in the $NHLF$ case, see the insert in Fig.1.

In Fig.3 we report the density dependence of neutron (bottom) and proton (top) effective masses for various asymmetries ($I = (N - Z)/A$) as predicted by $NLHF$ (solid lines) and $NLH + \rho + \delta$ (dashed lines). We remind that in the usual Hartree approximation this effect is associated with the scalar isovector δ meson. The Fock terms lead to a behaviour:

$$M^*_{n,p}(NLHF) = M^* \mp f^m_S(\rho_{Sn} - \rho_{Sp}) + \frac{b^2 + b_0^2}{16}\frac{d^2\Phi}{d\rho_S^2} \qquad (- \equiv n, + \equiv p), \tag{10}$$

where M^* is the nucleon effective mass in symmetric NM,

$$M^* = M - \Phi - (\frac{1}{2}f_V - \frac{1}{8}\frac{d\Phi}{d\rho_S} - \frac{1}{2}f_\rho)\rho_S + \frac{1}{16}(\rho_S^2 + \rho_B^2)\frac{d^2\Phi}{d\rho_S^2}. \tag{11}$$

and

$$f^m_S = \frac{f_V}{2} - \frac{1}{8}\frac{d\Phi}{d\rho_S} - \frac{f_\rho}{2} \tag{12}$$

with the isovector densities b, b_0 defined in Eq.(7). Since the coefficient f^m_S is positive we get an effect very similar to what expected from the contribution of the δ meson [18], dashed lines in Fig.3. The splitting of proton and neutron effective masses influences also the density behaviour of the symmetry energy (Fig.2) and is responsible for its rapid increase at high density. On this point we would like to make two more remarks:

i) The splitting is quite small around normal density, so it can be neglected in finite nuclei. This difference in n and p effective masses can however be relevant for transport properties of asymmetric, dense NM that can be formed in intermediate energy reactions with radioactive beams, naturally apart neutron star properties;

ii) The proton effective masses are systematically above the neutron ones. This trend, also in agreement with δ calculations [18,20], is just the opposite of what expected from Brueckner-Hartree-Fock calculations with realistic NN potentials [21]. Although relativistic and non-relativistic effective masses cannot be directly compared, [9], a clarification is needed, apart the possibility of an experimental check as said before.

In conclusion we have shown the evaluation, *in a non-perturbative scheme*, of Fock term contributions in a non linear effective field theory for asymmetric nuclear matter. Very reasonable and stimulating results for isospin effects on the nuclear EOS are obtained just from such minimal explicit many-body effects.

We stress again that at variance with non relativistic effective forces all the RMF models give a *stiff* potential symmetry term, i.e. more repulsive with increasing baryon density. Moreover the density dependence of the isovector couplings due to the Fock contributions leads to a new interesting effect (see the $NLHF$ curve in the insert of Fig.1), a *softening* of the behaviour at sub-nuclear densities. We expect to see related dynamical effects in heavy ion collisions at intermediate energies in fragmentation events [22] and collective flows [23,24].

Of course a similar analysis can be performed for spin effects. Moreover a transport equation can be consistently derived to be used for the study of dynamical evolution of nuclear matter far from normal conditions.

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FIGURE CAPTIONS

Figure 1 - Symmetry energy per nucleon vs. baryon density (in units of saturation density). Solid line: $NLHF$ results. Dashed line: Hartree results with ρ and δ mesons (NLH). Long-dashed: Hartree results with only ρ meson. In the insert we magnify the $NLHF - NLH$ comparison in the region below saturation density.

Figure 2 - Density dependence of f'_S and f'_V. Each curve is normalized to the value at saturation density.

Figure 3 - Proton (top curves) and neutron (bottom curves) effective masses vs. baryon density for various charge asymmetries. $I = \frac{N-Z}{N+Z} = 0:$ long dashed line. $I = 0.8:$ solid lines $NLHF$; dashed lines NLH.